

# Dirac gauginos, R symmetry and the 125 GeV Higgs

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## ABSTRACT:

We study a supersymmetric scenario with a quasi exact R-symmetry in light of the discovery of a Higgs resonance with a mass of 125 GeV. In such a framework, the additional adjoint superfields, needed to give Dirac masses to the gauginos, contribute both to the Higgs mass and to electroweak precision observables. We analyze the interplay between the two aspects, finding regions in parameter space in which the contributions to the precision observables are under control and a 125 GeV Higgs boson can be accommodated. We estimate the fine-tuning of the model finding regions of the parameter space still unexplored by the LHC with a fine-tuning considerably improved with respect to the minimal supersymmetric scenario. In particular, sizable non-holomorphic (non-supersoft) adjoints masses are required to reduce the fine-tuning.

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## 1 Introduction

The discovery of a 125 GeV particle closely resembling the Standard Model (SM) Higgs [1, 2] may represent a challenge for Supersymmetry (SUSY). Indeed, at least in its minimal version, large loop contributions are needed to raise the mass of the lightest Higgs boson to the observed value, the most relevant ones coming from the stop system. This points toward very heavy stops, and/or large left-right stop mixing.

While this is perfectly consistent with the non observation of any superpartner at the LHC, it is widely believed to be at odds with the concept of naturalness, which requires light stops with small left-right mixing. Needless to say, after the first LHC run and the Higgs discovery, understanding whether the concept of naturalness as it stands is or not a principle followed by nature has become of the utmost importance.

If we insist on naturalness, we need to consider alternatives to the Minimal Supersymmetric Standard Model (MSSM). An interesting possibility is given by models with Dirac gauginos, which have relaxed naturalness bounds on the gluino mass. This is most welcome, since being the gluino the most constrained particle after the first LHC run, a relaxed naturalness bound on its mass gives less tension with data. The mechanism behind the improved naturalness is the generation of Dirac gaugino masses through supersoft operators, which give only finite contributions to scalar masses [3]. Models with Dirac gauginos are also interesting from a purely phenomenological point

of view: first of all, squark pair production is suppressed at the LHC due to the absence of Majorana mass insertions, resulting in less stringent bounds on the gluino mass [4]. Moreover, Dirac gaugino masses are compatible with the presence of a global  $U(1)_R$  symmetry, which would be otherwise broken by the Majorana mass. The R-symmetry can be used as an alternative to  $R$  parity to forbid operators leading to proton decay [5, 6], but has far richer consequences. Indeed, the absence of  $A$  terms, the  $\mu$  term and Majorana gaugino masses has a drastic beneficial effect on the SUSY flavor problem [7].

A peculiar aspect of R-symmetric models is the Higgs sector particle content. Models have been proposed in the literature with four Higgs doublets [7], two Higgs doublets in which the role of the down type Higgs is played by one of the lepton doublets [8], one up type Higgs doublet [9] or even with no Higgs doublets at all, with the role of the Higgs being played by one of the slepton doublets [10].

As already pointed out, naturalness is among the reasons motivating the study of models with Dirac gauginos. However, a solid and complete statement about naturalness cannot be done without a full analysis of how a 125 GeV Higgs mass is obtained within this framework. The situation has been partially studied in [11], where however the R-symmetric case was not considered. This case is going to be the focus of this paper. As we will explain, respecting the R-symmetry in the Higgs sector changes dramatically how the lightest Higgs mass is raised up to 125 GeV. Indeed, while in [11] this is achieved through an NMSSM-like tree level enhancement of the Higgs mass, here this possibility is forbidden by the R-symmetry. However, it turns out that the extra matter necessary to respect the R-symmetry, *i.e.* the adjoint scalars and the inert doublets, can provide radiative corrections comparable to the stop one, giving a 125 GeV Higgs with a few percent level fine-tuning.

## 2 Electroweak symmetry breaking in R-symmetric models

As already explained, preserving the R-symmetry typically requires an enlarged Higgs sector. For definiteness, we will present the Lagrangian for the four Higgs doublet model [7], in which the two doublets with  $R$  charge 0,  $H_u$  and  $H_d$ , acquire a vev while the two with  $R$  charge 2,  $R_u$  and  $R_d$ , are inert doublets. Another, more economical, possibility is to have the sneutrino as the down type Higgs so that just two doublets,  $H_u$  and  $R_d$  are needed.<sup>1</sup> For reasons which will become clear later on, we will focus on the large  $\tan\beta$  limit (were  $\tan\beta \gtrsim 10$ ) in which most of electroweak symmetry breaking

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<sup>1</sup>It is also possible to have an even more economical Higgs sector [10] where the sneutrino gives mass to the up type fermions via SUSY breaking Yukawa couplings. However, in this case the Higgs quartic is generated by SUSY breaking as well.

is through  $H_u$ , and in which the extra Higgs states decouple from the electroweak symmetry breaking sector. We expect that the various models will give similar results.

The superpotential of the model is given by:

$$\begin{aligned}
W &= W_{Yukawa} + W_{Higgs} \\
W_{Yukawa} &= H_u Q Y_u u^c + H_d Q Y_d d^c + H_d L Y_e e^c, \\
W_{higgs} &= \sqrt{2} \lambda_T^u H_u T R_d + \sqrt{2} \lambda_T^d R_u T H_d + \lambda_S^u H_u S R_d + \lambda_S^d R_u S H_d \\
&\quad + \mu_u H_u R_d + \mu_d R_u H_d.
\end{aligned} \tag{2.1}$$

We write the triplet superfield normalized as

$$T = \frac{1}{\sqrt{2}} \begin{pmatrix} T^0 & \sqrt{2} T^+ \\ \sqrt{2} T^- & -T^0 \end{pmatrix}, \tag{2.2}$$

so that the kinetic terms for the (complex) triplet components are automatically canonically normalized; the factor  $\sqrt{2}$  in front of  $\lambda_T^i$  is chosen such that  $W \supset \lambda_T H_u^0 T^0 R_d^0$ .

The R-symmetry allows the gaugino fields  $\lambda_i$  to pair up with the fermionic components of the adjoint superfields,  $\psi_i$ , through soft SUSY breaking Dirac masses

$$\mathcal{L}_D = M_{\tilde{B}} \lambda_{\tilde{B}} \psi_{\tilde{B}} + M_{\tilde{W}} \lambda_{\tilde{W}}^a \psi_{\tilde{W}}^a + M_{\tilde{g}} \lambda_{\tilde{g}} \psi_{\tilde{g}} + h.c. \tag{2.3}$$

Moreover, the soft SUSY breaking scalar terms read

$$\begin{aligned}
V_{\text{soft}}^{EW} &= \tilde{Q}^\dagger m_{\tilde{Q}}^2 \tilde{Q} + \tilde{u}^\dagger m_{\tilde{u}}^2 \tilde{u} + \tilde{d}^\dagger m_{\tilde{d}}^2 \tilde{d} + \tilde{L}^\dagger m_{\tilde{L}}^2 \tilde{L} + \tilde{e}^\dagger m_{\tilde{e}}^2 \tilde{e} + B_\mu H_u H_d \\
&\quad + m_{H_u}^2 |H_u|^2 + m_{H_d}^2 |H_d|^2 + m_{R_u}^2 |R_u|^2 + m_{R_d}^2 |R_d|^2 + m_s^2 |S|^2 + m_T^2 T^{a\dagger} T^a \\
&\quad + t_S S + B_S S^2 + \frac{1}{3} A_S S^3 + B_T T^{a\dagger} T^a \\
&\quad + A_{ST} S T^2 + A_{SH} S H_u H_d + A_{TH} H_u T H_d + h.c.
\end{aligned} \tag{2.4}$$

We notice that the R-symmetry forbids all the  $A$ -terms except for those written above, which together with  $t_S$  we will assume to be negligible for simplicity.<sup>2</sup>

Let us now comment on the soft breaking terms in the adjoint sector. As already explained in the introduction, Dirac gaugino masses are generated by supersoft operators and give finite contributions to the scalar masses. This property has the beneficial effect of relaxing the gluino naturalness bound, reducing the tension with the direct searches [4]. However not all possible R-invariant terms that can be constructed out of

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<sup>2</sup>Note that obtaining  $t_S$  small could require some fine-tuning which we do not quantify here. However, we have checked that allowing  $t_S$  to be of order  $v^3$  does not qualitatively change our findings.

the adjoint superfields are supersoft: indeed the non-holomorphic adjoints masses for the singlet  $m_S^2$ , the triplet  $m_T^2$ , and the octet  $m_O^2$  contribute at the two-loop level to the  $\beta$  functions for the scalar masses, pushing down their values at low energy. In particular, a too large octet scalar mass can eventually induce tachyonic squark masses, causing charge and color breaking at the weak scale [12, 13]. Furthermore, it is also important for these three terms (Dirac gaugino masses, holomorphic and non-holomorphic adjoints scalar masses) to be of the same order, to avoid tachyons already at tree level. It turns out, however, that realizing this spectrum in a UV complete model is quite challenging. This resembles the  $\mu - B_\mu$  problem in gauge mediation, and it leads to a source of fine-tuning estimated in [13] to be of order of 0.1%. In what follows we will discuss a generic case where also non-holomorphic masses are present, assuming that the mass hierarchy among the adjoint soft terms is such to ensure color and charge conservation at the weak scale. In fact, as we will discuss in more detail in Sec. 4, the presence of non-holomorphic masses is instrumental in order to get a 125 GeV Higgs improving at the same time the fine-tuning with respect to the MSSM.

The total scalar potential is

$$V^{EW} = V_F^{EW} + V_D^{EW} + V_{\text{soft}}^{EW} ,$$

$$V_F^{EW} = \sum_i \left| \frac{\partial W}{\partial \phi_i} \right|^2 , \quad V_D^{EW} = \frac{1}{2} \sum_{a=1}^3 (D_2^a)^2 + \frac{1}{2} D_Y^2 , \quad (2.5)$$

with  $W$  defined in Eq. (2.1) and  $V_{\text{soft}}^{EW}$  given in Eq. (2.4). The presence of additional chiral superfields charged under  $SU(2)_L \times U(1)_Y$  modifies the expression for the D-terms:

$$D_2^a = g \left( H_u^\dagger \tau^a H_u + H_d^\dagger \tau^a H_d + R_u^\dagger \tau^a R_u + R_d^\dagger \tau^a R_d + \vec{T}^\dagger \lambda^a \vec{T} \right) + \sqrt{2} M_{\tilde{W}} \left( \vec{T}^a + \vec{T}^{\dagger a} \right) ,$$

$$D_Y = \frac{g'}{2} \left( H_u^\dagger H_u + R_u^\dagger R_u - H_d^\dagger H_d - R_d^\dagger R_d \right) + \sqrt{2} M_{\tilde{B}} (S + S^\dagger) , \quad (2.6)$$

where  $M_{\tilde{B}}$  and  $M_{\tilde{W}}$  are the Dirac Bino and Wino masses,  $\tau^a$  and  $\lambda^a$  are the two and three dimensional  $SU(2)$  generators respectively, while  $\vec{T}^a = \sqrt{2} \text{tr}(\tau^a T) = \left\{ \frac{T^+ + T^-}{\sqrt{2}}, \frac{T^- - T^+}{\sqrt{2}i}, T^0 \right\}$ .

Writing the neutral fields as

$$H_{u,d}^0 = \frac{h_{u,d} + ia_{u,d}}{\sqrt{2}}, \quad R_{u,d}^0 = \frac{r_{u,d} + ia_{r_{u,d}}}{\sqrt{2}}, \quad T^0 = \frac{t + ia_t}{\sqrt{2}}, \quad S = \frac{s + ia_s}{\sqrt{2}}, \quad (2.7)$$

the scalar potential for the CP even components reads:

$$\begin{aligned}
V = & \frac{1}{2} \left[ (m_{H_u}^2 + \mu^2) h_u^2 + (m_{R_u}^2 + \mu^2) r_u^2 + (m_{H_d}^2 + \mu^2) h_d^2 + (m_{R_d}^2 + \mu^2) r_d^2 \right. \\
& - 2B_\mu h_u h_d + (4M_{\tilde{B}}^2 + m_S^2 + 2B_S) s^2 + (4M_{\tilde{W}}^2 + m_T^2 + 2B_T) t^2 \Big] \\
& + \frac{1}{2} \left[ \sqrt{2}\mu (\lambda_S s + \lambda_T t) (h_u^2 + h_d^2 + r_u^2 + r_d^2) \right. \\
& + (gM_{\tilde{W}} - g'M_{\tilde{B}}) (h_d^2 + r_d^2 - h_u^2 - r_u^2) \Big] + \frac{1}{32} (g^2 + g'^2) \left[ (h_u^2 - h_d^2)^2 + (r_u^2 - r_d^2)^2 \right] \\
& + \frac{g^2 + g'^2}{16} (h_u^2 r_u^2 + h_d^2 r_d^2) + \left( \frac{\lambda_S^2 + \lambda_T^2}{4} - \frac{g^2 + g'^2}{16} \right) (h_u^2 r_d^2 + h_d^2 r_u^2) \\
& + \frac{\lambda_S \lambda_T}{2} (h_u^2 + h_d^2 + r_u^2 + r_d^2) st + \frac{\lambda_T^2 t^2 + \lambda_S^2 s^2}{4} (h_u^2 + h_d^2 + r_u^2 + r_d^2), \tag{2.8}
\end{aligned}$$

where we have assumed for simplicity  $\lambda_S^u = \lambda_S^d = \lambda_S$ ,  $\lambda_T^u = \lambda_T^d = \lambda_T$ , and  $\mu_u = \mu_d = \mu$  and set  $A_{ST} = A_{SH} = A_{TH} = 0$ . The minimization conditions for this potential are written in the appendix. The triplet acquires a vev which is constrained by EWPM to be  $|v_T| \lesssim 3$  GeV. We will discuss more precisely the bounds from EWPM in Sec. 3.

Inspecting the various contributions, we notice that the D-terms produce the usual MSSM quartic. However, the Dirac gaugino masses contribute to reduce the tree level Higgs mass with respect to the MSSM. Indeed,  $V_D$  contains trilinear interactions between the active Higgs fields (those participating in EWSB) and the scalar adjoints,

$$V_D \supset \frac{1}{2} (-gM_{\tilde{W}}t + g'M_{\tilde{B}}s) h_u^2 - \frac{1}{2} (-gM_{\tilde{W}}t + g'M_{\tilde{B}}s) h_d^2, \tag{2.9}$$

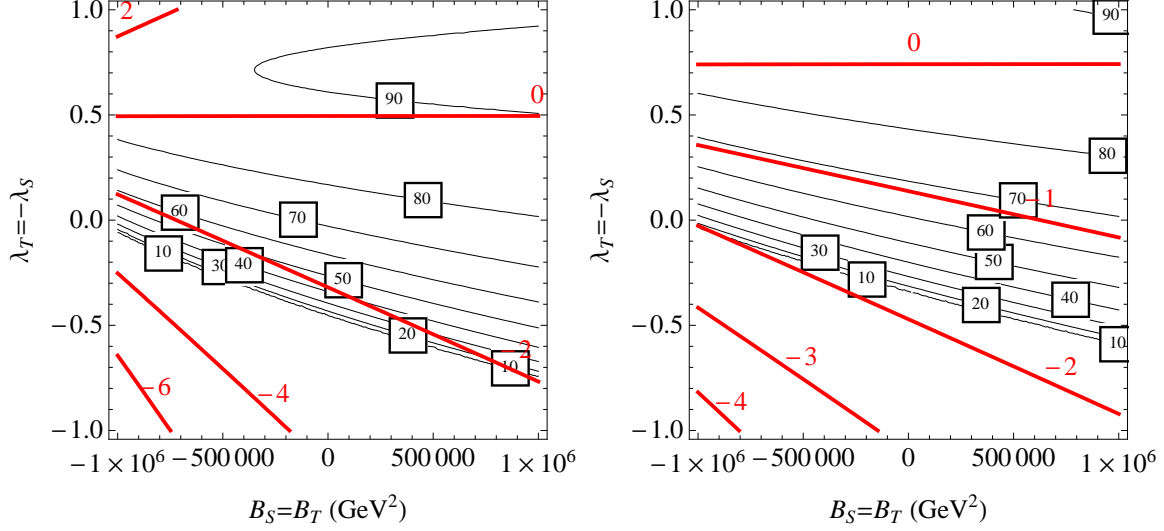
which after EWSB push down the lightest eigenvalue due to mixing.

In addition, the R-symmetry forces the active Higgs fields to couple only with the inert doublets (those that do not get vevs) and not among themselves, so that any NMSSM-like quartic term  $\lambda_{S,T}^2 h_u^2 h_d^2$  is forbidden. As a consequence, the MSSM tree level upper bound  $(m_h)_{tree}^2 \leq m_Z^2 \cos^2 2\beta$  applies, and the lightest scalar mass is maximized in the large  $\tan \beta$  regime. The situation is different when the R-symmetry is broken in the Higgs sector. In this case  $W \supset \lambda_T H_u T H_d + \lambda_S H_u S H_d$  and in the low  $\tan \beta$  regime the usual NMSSM-like tree level enhancement is recovered [11].

A more complete discussion of the tree level scalar masses will be presented in Sec. 2.1, where in order to maximize the lightest eigenvalue we will focus on the large  $\tan \beta$  regime.<sup>3</sup> In Sec. 2.2 we will instead study the loop corrections to the Higgs boson mass.

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<sup>3</sup>However, we have checked that it is easy to deform our benchmark points to obtain examples with moderate  $\tan \beta$  ( $\sim 10$ ) without affecting our conclusions.



**Figure 1.** Tree level Higgs boson mass in GeV (black lines) and singlet vev  $v_S$  (red lines) as a function of  $B_T = B_S$  and  $\lambda_T = -\lambda_S$ . Left:  $M_{\tilde{W}} = M_{\tilde{B}} = 600$  GeV,  $m_T = m_S = 1500$  GeV and  $\mu = 300$  GeV. Right:  $M_{\tilde{W}} = M_{\tilde{B}} = 900$  GeV,  $m_T = m_S = 1500$  GeV and  $\mu = 300$  GeV.

## 2.1 Tree level Higgs mass

We already pointed out that Dirac gaugino masses constitute an irreducible source of mixing between active Higgs fields and adjoint scalars. This push-down effect may in part be kept under control by the supersymmetric couplings  $\lambda_T$ ,  $\lambda_S$  and by the  $\mu$  term. This is evident looking at the off diagonal elements of the mass matrix for CP even scalars (see Appendix A for the complete expressions):

$$\begin{aligned} m_{h_u, t}^2 &= v(-\sqrt{2}gM_{\tilde{W}} + 2\lambda_T(\lambda_S v_S + \lambda_T v_T + \mu)), \\ m_{h_u, s}^2 &= v(+\sqrt{2}g'M_{\tilde{B}} + 2\lambda_S(\lambda_S v_S + \lambda_T v_T + \mu)). \end{aligned} \quad (2.10)$$

Anticipating that  $\lambda$  couplings of order one are helpful to increase the Higgs boson mass at loop level when the stops are not too heavy, and insisting on relatively small  $\mu$  values as suggested by naturalness (see Eq. (4.3)), we see that the terms in Eq. (2.10) can be kept under control for small singlet and triplet vevs. This arises from a partial cancellation between the first and last terms. Since by field redefinitions we can always choose  $g > 0$  and  $M_{\tilde{B}}, M_{\tilde{W}}, \mu > 0$ , this then implies that  $\lambda_T > 0$  and  $\lambda_S < 0$  are preferred to obtain smaller  $m_{h_u, t}^2$  and  $m_{h_u, s}^2$ . This is confirmed in Fig. 1, where we show the tree level Higgs boson mass, together with the singlet vev  $v_S$ , as a function

of  $B_T = B_S$  and  $\lambda_T = -\lambda_S$ . In both cases, it is possible to get  $m_h \simeq m_Z$  for small and positive  $v_S$ .

Going back to the mixing between  $h_u$ ,  $s$  and  $t$  and the related mass reduction, a simple formula can be obtained in the limit of small  $v_T$ ,  $v_S$  and large hierarchy between Dirac gaugino masses and non-holomorphic adjoint masses,  $M_D \ll m_{adj}$ . For  $\tan \beta \gg 1$ , the lightest tree level mass is:

$$(m_h^2)_{tree} \simeq m_Z^2 - v^2 \frac{(-\sqrt{2}gM_{\tilde{W}} + 2\lambda_T\mu)^2}{m_{T_R}^2} - v^2 \frac{(\sqrt{2}g'M_{\tilde{B}} + 2\lambda_S\mu)^2}{m_{S_R}^2}, \quad (2.11)$$

where  $m_{T_R}^2 = 4M_{\tilde{W}}^2 + m_T^2 + 2B_T$  and  $m_{S_R}^2 = 4M_{\tilde{B}}^2 + m_S^2 + 2B_S$  are the masses of the real adjoint scalars before EWSB. Let us stress that the presence of supersymmetric couplings, as well as holomorphic and non-holomorphic masses for the adjoint scalars, improves the situation with respect to [3], where the quartic coupling vanishes for decoupled adjoint scalars (see also [11]).

## 2.2 Radiative corrections to the Higgs mass

We compute now the 1-loop corrected Higgs mass using the Coleman-Weinberg potential:

$$V_{Higgs}^{CW} = \frac{1}{64\pi^2} \left[ \sum_i (-1)^{2J_i+1} (2J_i + 1) m_i^4 \left( \log \frac{m_i^2}{Q^2} - \frac{3}{2} \right) \right]. \quad (2.12)$$

The sum is to be taken over all the states coupled to the Higgs, with  $m_i^2$ 's the field dependent masses. To obtain analytic expressions for the loop corrections, we expand the field dependent masses in powers of  $h_u$ , setting to zero the singlet and triplet backgrounds (we know from the previous section that  $v_S$  must be small in order for the tree level Higgs mass not to be too different from  $m_Z$ , while  $v_T$  must be small to fulfill the precision measurement constraints). We do not present here the full analytical expressions, since they are lengthy and not particularly transparent. Simple expressions can be obtained for  $M_D \ll m_{adj}$  or  $m_{adj} \ll M_D$ , where  $M_D$  and  $m_{adj}$  are common mass scales for Dirac gauginos and adjoint scalars, respectively.

**Region 1.** When the scalar CP even and CP odd masses are significantly larger than the gaugino masses,  $\mu \ll M_D \ll m_{adj}$ , we have the following contribution to the



Higgs quartic coupling:

$$\begin{aligned}
V_{Higgs}^{CW} \supset & \frac{1}{4} \left[ \frac{5\lambda_T^4}{32\pi^2} \log \frac{m_T^2}{Q^2} + \frac{\lambda_S^4 + 2\lambda_T^2\lambda_S^2}{32\pi^2} \log \frac{m_S^2}{Q^2} \right. \\
& + \frac{5\lambda_T^4 + 2\lambda_T^2\lambda_S^2 + \lambda_S^4}{32\pi^2} \log \frac{m_{R_d}^2}{Q^2} - \frac{\lambda_T^2\lambda_S^2}{16\pi^2} \left. \right] h_u^4 \\
& - \frac{1}{4} \left[ \frac{5\lambda_T^4}{16\pi^2} \log \frac{M_{\tilde{W}}^2}{Q^2} + \frac{\lambda_S^4}{16\pi^2} \log \frac{M_{\tilde{B}}^2}{Q^2} + \frac{\lambda_S^2\lambda_T^2}{16\pi^2} \right] h_u^4 \quad (2.13)
\end{aligned}$$

where  $Q$  is the renormalization scale and the first two lines show the scalar contribution while the third one shows the fermionic one. We checked that this expression is still a good approximation in the more interesting limit where milder hierarchies among the masses hold. However, in what follows we will use the exact expressions to compute the Higgs boson mass. A particularly simple expression can be obtained in the limit  $m_{R_d}^2 \simeq m_T^2 \simeq m_S^2$ . The Higgs quartic is then

$$V_{Higgs}^{CW} \supset \frac{1}{4} \left[ \frac{5\lambda_T^4}{32\pi^2} \log \frac{m_T^2}{M_{\tilde{W}}^2} + \frac{\lambda_S^4}{32\pi^2} \log \frac{m_S^2}{M_{\tilde{B}}^2} - \frac{\lambda_T^2\lambda_S^2}{8\pi^2} \right] h_u^4, \quad (2.14)$$

so that a relevant positive contribution to the quartic can be obtained for a large enough ratio  $m_{adj}/M_D$ .

**Region 2.** In the opposite limit,  $m_{adj} \ll M_D$ , the one-loop contribution to the Higgs quartic is :

$$\begin{aligned}
V_{Higgs}^{CW} \supset & \frac{1}{4} \left[ \frac{5\lambda_T^4}{32\pi^2} \log \frac{M_{\tilde{W}}^2}{Q^2} + \frac{\lambda_S^4 + 2\lambda_T^2\lambda_S^2}{32\pi^2} \log \frac{M_{\tilde{B}}^2}{Q^2} \right] h_u^4 \\
& - \frac{1}{4} \left[ \frac{5\lambda_T^4}{16\pi^2} \log \frac{M_{\tilde{W}}^2}{Q^2} + \frac{\lambda_S^4}{16\pi^2} \log \frac{M_{\tilde{B}}^2}{Q^2} - \frac{\lambda_T^2\lambda_S^2}{8\pi^2} \right] h_u^4 \quad (2.15)
\end{aligned}$$

where we have also assumed  $m_{R_d} \ll M_D$ . The first line shows the scalar contribution, the second line the fermionic one.

Putting all together, we end up with

$$V_{Higgs}^{CW} \supset \frac{1}{4} \left[ -\frac{5\lambda_T^4}{32\pi^2} \log \frac{M_{\tilde{W}}^2}{Q^2} - \frac{\lambda_S^4 - 2\lambda_T^2\lambda_S^2}{32\pi^2} \log \frac{M_{\tilde{B}}^2}{Q^2} + \frac{\lambda_T^2\lambda_S^2}{8\pi^2} \right] h_u^4 \quad (2.16)$$

so that we expect this contribution to be always negative. Let us notice that this region corresponds to the pure supersoft spectrum, where indeed the non-holomorphic scalar masses are negligible and  $M_D^2 \gtrsim B$  in order to avoid problems

with tachyonic masses. Furthermore,  $m_{R_d}^2$  is given by the gaugino induced one-loop correction [3]:

$$m_{R_d}^2 = \frac{\alpha_2}{\pi} M_W^2 \log \frac{4M_W^2 - 2B_T}{M_{\tilde{W}^2}}, \quad (2.17)$$

with an inert doublet therefore too light to give any significant boost to the Higgs mass.

For comparison, the well known stop contribution is given by [14]

$$V_{Higgs}^{CW} \supset \frac{1}{4} \left[ \frac{3}{16\pi^2} y_t^2 \left( y_t^2 - \frac{m_Z^2}{2v^2} \right) \log \frac{M^2}{m_t^2} + \frac{3y_t^4}{(16\pi^2)^2} \left( \frac{3}{2} y_t^2 - 32\pi\alpha_3(m_t) \right) \log^2 \frac{M^2}{m_t^2} \right] h_u^4, \quad (2.18)$$

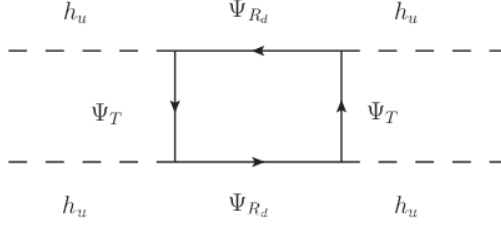
where  $y_t$  is the top Yukawa coupling,  $\alpha_3$  the strong coupling constant and  $M$  is a common soft SUSY breaking stop mass scale. We show also the two-loop contribution, since the term proportional to the strong gauge coupling may reduce in a significant way the Higgs quartic.

The simplified expressions, Eqs. (2.13), (2.14), suggest that for  $\lambda_T \simeq \lambda_S \simeq y_t$  the new states may give a contribution comparable to the stop one, depending on the mass hierarchy. In this sense, they can be regarded as “additional stops”, which in principle can allow for a collective loop enhancement of the Higgs quartic. Moreover, since both the triplet and the singlet are uncolored, we do not expect the two loop terms proportional to the gauge couplings to give a reduction analogous to the one proportional to  $\alpha_3$  in the stop sector, making more effective the loop boost achieved through these states. Whether or not this scenario will allow to obtain a 125 GeV Higgs with less fine-tuning than in the MSSM will be studied in detail in Sec. 4.

### 3 Electroweak Precision Measurements

Getting a significant help from the triplet and the singlet to raise the Higgs mass through radiative corrections requires an appreciable value for the couplings  $\lambda_T$  and/or  $\lambda_S$ . However, the same couplings contribute to the  $T$  parameter at loop level. Therefore there exist a potential tension between generating a large Higgs mass and electroweak precision data. Besides the loop-level corrections to  $T$  there is already at tree level a dangerous effect due to the vev for the triplet  $T^0$ , which can lead to a large contribution to the  $\hat{T}$  parameter (with the standard  $T = \hat{T}/\alpha$ ):

$$\hat{T} = 4 \frac{v_T^2}{v^2}, \quad (3.1)$$



**Figure 2.** Box diagram that contributes to the  $T$  parameter.

which constrain the triplet vev to be  $|v_T| \lesssim 3$  GeV, where

$$v_T = \frac{\sqrt{2}gM_{\tilde{W}} - 2\lambda_S\lambda_T v_S - 2\lambda_T\mu}{4B_T + 8M_{\tilde{W}}^2 + 2m_T^2 + 2\lambda_T^2 v^2} v^2. \quad (3.2)$$

It can be minimized by taking  $m_T$  large, or otherwise arranging for the numerator to be small. Besides the tree level contributions, there are contributions coming from loop of superpartners. A detailed study of all these contributions will be presented in [15], but the dominant effect comes from contributions to  $\hat{T}$  from loops involving the fermionic part of the superfield  $H_u$ ,  $T$ ,  $R_d$  and  $S$ . Integrating them out at loop level, through the diagram of Fig. 2 lead to the higher-dimension operator associated with  $\hat{T}$ :

$$\frac{|H_u^\dagger D_\mu H_u|^2}{\Lambda^2}, \quad (3.3)$$

with a coefficient proportional to  $\lambda_T^4$ . Thus the same coupling which can help to make the Higgs heavier will also lead to too large contributions to  $T$ . To estimate the region of parameter space excluded by electroweak precision data we compute the  $\hat{T}$  parameter due to  $v_T$ , and from loops of the scalar and fermionic sector of  $H_u$ ,  $T$ ,  $R_d$  and  $S$ . Imposing  $T < 0.2$  [16], we find that the fermion loops will force  $\lambda_T \lesssim 1$ , whenever  $M_D \lesssim 1$  TeV.

## 4 125 GeV Higgs boson and fine-tuning

We are now in the position to analyze the region in parameter space in which not only a 125 GeV Higgs boson mass can be obtained, but also the contributions to the  $T$  parameter can be kept under control. Before doing so we turn to the fine-tuning issue. Taking as measure [17]

$$\Delta = \max_{a_i} \frac{\partial \log m_h^2}{\partial \log a_i}, \quad (4.1)$$

where  $a_i$  are the lagrangian parameters of our model. We find that the fine-tuning is dominated by the mass parameters:  $m_T^2, m_S^2, m_{R_d}^2, m_{stop}^2$  and  $\mu^2$ . Using for the RGE's the results available in the literature [18], we get the following  $\Delta$ -dependent upper bounds on the soft SUSY breaking masses:

$$\begin{aligned}
m_{\tilde{Q}, \tilde{u}} &\lesssim 600 \text{ GeV} \frac{m_h}{125 \text{ GeV}} \sqrt{\frac{3}{\log \Lambda / 1 \text{ TeV}}} \sqrt{\frac{\Delta}{5}} , \\
m_T &\lesssim \frac{600 \text{ GeV}}{|\lambda_T|} \frac{m_h}{125 \text{ GeV}} \sqrt{\frac{3}{\log \Lambda / 1 \text{ TeV}}} \sqrt{\frac{\Delta}{5}} , \\
m_S &\lesssim \frac{1000 \text{ GeV}}{|\lambda_S|} \frac{m_h}{125 \text{ GeV}} \sqrt{\frac{3}{\log \Lambda / 1 \text{ TeV}}} \sqrt{\frac{\Delta}{5}} , \\
m_{R_d} &\lesssim \frac{1000 \text{ GeV}}{\sqrt{6\lambda_T^2 + 2\lambda_S^2}} \frac{m_h}{125 \text{ GeV}} \sqrt{\frac{3}{\log \Lambda / 1 \text{ TeV}}} \sqrt{\frac{\Delta}{5}} .
\end{aligned} \tag{4.2}$$

In addition, there is also a tree level upper bound on the  $\mu$  parameter:

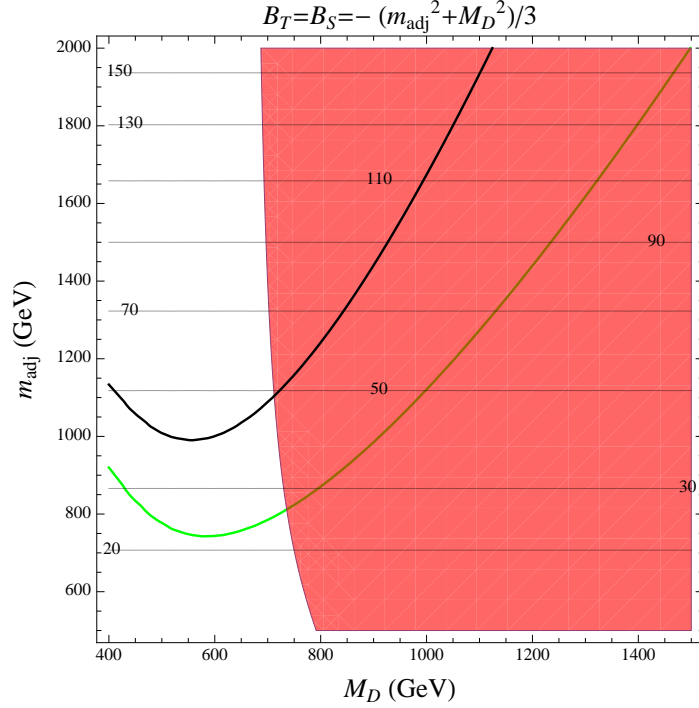
$$\mu \lesssim 200 \text{ GeV} \sqrt{\frac{\Delta}{5}} . \tag{4.3}$$

In contrast to what happens in models with Majorana gaugino masses, there is no dependence on the Dirac gaugino masses in the RGE's for the  $m_{H_u}^2$  parameter.<sup>4</sup> We also see from Eq. (4.2) that for  $\lambda_T \sim \lambda_S \sim 1$ , the fine-tuning is driven by  $m_{R_d}$ . Moreover, since for  $\lambda_T \sim 1$  the fine-tuning due to  $m_{stop}$  is of the same order than the one due to  $m_{adj}$ , there is no worsening of the fine-tuning when  $m_{\tilde{t}} \sim m_{adj}$ . Our main results are shown in Figs. 3 and 4. In Fig. 3 we present the results as a function of  $M_D = M_{\tilde{W}} = M_{\tilde{B}}$  and  $m_{adj} = m_T = m_S = m_{R_d}$ , fixing also  $\lambda_T = 1 = -\lambda_S$ . The solid lines corresponds to  $m_h = 125 \text{ GeV}$ , the blue region is the one in which there are unwanted tachyon, while the red region is allowed at 95% C.L. by EWPM. We also show the fine-tuning parameter  $\Delta$  (thin black lines), Eq. (4.1), where we fix  $\Lambda$ , the scale at which the soft terms are generated, at 20 TeV.<sup>5</sup> We use two different values for the stop mass:  $m_{\tilde{t}} = 300 \text{ GeV}$  (upper curve) and  $m_{\tilde{t}} = m_{adj}$  (lower curve). The first case refers

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<sup>4</sup>There is nonetheless a dependence through the finite one-loop contribution analogous to Eq. (2.17); we checked however that the tuning due to this contribution is never the dominant one in the interesting regions.

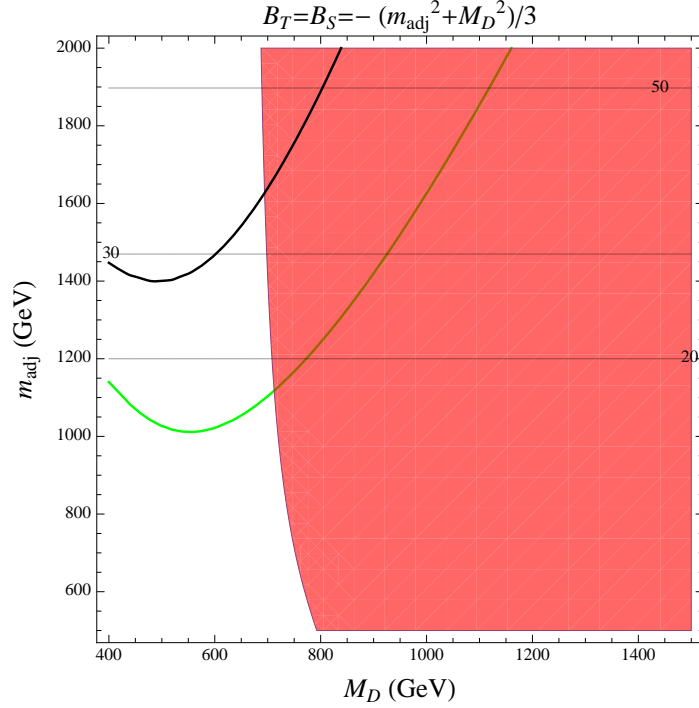
<sup>5</sup>In this work we assume that the required soft parameters can be obtained naturally with the appropriate values at this scale. However, as mentioned previously, this is a somewhat non-trivial task and could be the source of additional fine-tuning [12].



**Figure 3.** Higgs boson mass  $m_h = 125$  GeV (thick lines) and fine-tuning parameter (thin lines), as a function of  $M_D = M_{\tilde{W}} = M_{\tilde{B}}$  and  $m_{adj} = m_T = m_S = m_{R_d}$ , and  $B_T = B_S = -\frac{1}{3}(m_{Adj}^2 + M_D^2)$ . We fix  $\lambda_T = 1 = -\lambda_S$ ,  $\mu = 300$  GeV and the common stop mass to  $m_{stop} = 300$  GeV for the upper (black) curve and  $m_{stop} = m_{adj}$  for the lower (green) curve. Red region: allowed at 95% C.L. by EWPM ( $T < 0.2.$ )

to a situation with a very light stop<sup>6</sup> in which the main boost to the Higgs quartic comes from the adjoint and inert fields. However, since the fine-tuning is dominated by  $m_{R_d}$ , a heavier stop is most welcome, since it gives a larger loop contribution to the Higgs quartic without worsening the fine-tuning (again, from Eq. 4.2 we see that the stops give a fine-tuning on the same order as  $m_{adj}$ ). This can be clearly seen from the green (lower) curve, where we choose  $m_{\tilde{t}} = m_{adj}$ . We see that for  $m_{\tilde{t}} \sim 700$  GeV, which is around the present direct searches exclusion limits [19, 20], there is an acceptable point with roughly  $\Delta \sim 30$ . For comparison, in the MSSM the minimal tuning needed to accommodate a 125 GeV Higgs is  $\Delta \gtrsim 100$  in the case of maximal stop mixing, or  $\Delta \gtrsim 300$  in the case of vanishing stop mixing [21].

<sup>6</sup> A detailed study of the LHC phenomenology of the model is outside the scope of the present work, therefore we assume  $m_{\tilde{t}} \sim 300$  GeV to be still allowed by the LHC either because the spectrum is very compressed, or because there are baryonic R-Parity violating couplings in the superpotential.

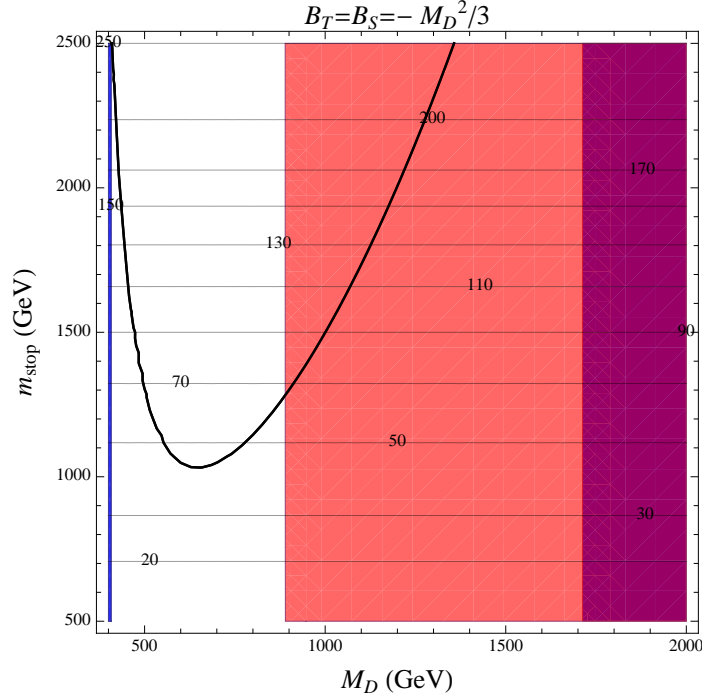


**Figure 4.** As in Fig. 3, but fixing  $m_{R_d} = m_{adj}/2$ .

Overall, compatibility between the Higgs boson mass and EWPM is achieved for  $M_D \gtrsim 700$  GeV,  $m_{adj} \gtrsim 800$  GeV, and for  $\mu > 300$  GeV; a lower value for  $\mu$  requires heavier gauginos and thus heavier scalars to be safe from the EWPM constraints. Moreover, the choice  $|\lambda_{S/T}| \sim 1$  is essential in order to achieve  $m_h = 125$  GeV with states close to the TeV scale. Lower values of the  $\lambda$  couplings would require either very heavy adjoint scalars or heavier stops.

In Fig. 4 we show the situation in which  $m_{R_d}$  is smaller than the adjoint masses, namely  $m_{R_d} = m_{adj}/2$ , again for the two cases  $m_{\tilde{t}} = 300$  GeV (upper black curve) and  $m_{\tilde{t}} = m_{adj}$  (lower green curve). We notice an overall improvement of the fine-tuning, up to  $\Delta \lesssim 20$ . This is due to the fact that the contribution to fine-tuning from  $m_{R_d}$ , which was in the previous situation the dominant one, is now reduced. The loss in Higgs mass can be compensated by a moderate increase of  $m_{adj}$ , resulting in a slightly more advantageous situation.

As a final comment, we consider the case of a pure supersoft spectrum, *i.e.* the case in which we set the non-holomorphic adjoint masses  $m_T$  and  $m_S$  to zero. With  $m_{R_d}^2$  given by (2.17), we easily see from Eq. (2.16) that the only relevant radiative correction comes from the stop sector, with the gaugino contributions decreasing the



**Figure 5.** Higgs boson mass  $m_h = 125$  GeV (solid line) and fine-tuning parameter (thin lines), as a function of  $M_D = M_{\tilde{W}} = M_{\tilde{B}}$  and  $m_{\tilde{t}} = m_{R_d}$ , with  $m_T = m_S = 0$  (supersoft limit), and  $B_T = B_S = -\frac{M_D^2}{3}$ . We fix  $\lambda_T = 1 = -\lambda_S$  and  $\mu = 300$  GeV. Blue region: spontaneous charge and/or CP breaking. Red region: allowed at 95% C.L. by EWPM ( $T < 0.2$ ). Purple region:  $m_{\tilde{l}_R} > 100$  GeV

Higgs mass. However, it is easy to imagine that the full UV completion of the model contains a sector which couples the Higgs multiplets to messenger fields. This can generate the required  $\mu$  and  $b_\mu$  terms, as well as extra contributions to the  $R_d$  mass which may then push up the Higgs mass even with moderately light stops. In any case, relying only on the loop corrections from  $R_d$  makes the boost to the Higgs quartic less efficient than in the non-supersoft case, and this makes the model less natural, since  $m_{R_d}^2$  is associated with a larger fine-tuning. There is also another important constraint to take into account in a pure supersoft spectrum: all the sfermions acquire mass via finite gaugino one loop contribution and are therefore predicted (modulo the running from the adjoint mass scale down to the weak scale) in term of  $M_D$ . This tends to make the sleptons, especially the right-handed ones which only have hypercharge gauge couplings, quite light. This is illustrated in Fig. 5, where we show once again, in the  $M_D$  vs  $M_{\tilde{t}}$  plane, the line corresponding to a 125 GeV Higgs together with the region

allowed by EWPM (in red). We also show the region where the slepton has a mass greater than 100 GeV (purple region), which correspond to the LEP bound. We see that the slepton constraint pushes all the masses to be very heavy, into a region with very large fine-tuning. This leads to the conclusion that sizable non-holomorphic adjoints masses are required to reduce the fine-tuning.

## 5 Conclusions

We are finally in an era in which experiments are directly exploring the electroweak scale, and will (at least in part) shed light on whether or not the electroweak scale is natural. The first LHC run has already provided us with some indications. The general message seems to be that our simplest natural models are by now tuned at the percent level, or worse. While it can turn out that this is the level of tuning of the EW scale, it may also be taken to motivate the search for more natural (although less minimal) models. One such possibility is the supersymmetric model with a quasi exact  $U(1)_R$  symmetry considered in this work. This kind of models are more natural with respect to the MSSM, since there is no gluino induced one-loop contribution to the squark masses. Moreover, the usual supersymmetric flavor problem is greatly ameliorated. The point on which we focus here is that the adjoint superfields needed to write Dirac gaugino masses may give relevant loop corrections to the Higgs boson mass: they act effectively as “additional stops”, at least in part of the parameter space. A possible drawback is that the very same couplings that help increasing the Higgs boson mass break custodial symmetry, potentially leading to large contributions to the electroweak precision measurements. Our main results are summarized in Figs. 3-4, in which we show the region in parameter space in which a 125 GeV Higgs mass can be obtained in a way compatible with EWPM. We also presented on the same plot the required fine-tuning. A first conclusion that can be drawn is that there are regions in which the fine-tuning is ameliorated with respect to the MSSM, roughly reduced at the same level as the NMSSM,  $\Delta \sim 20 - 30$  [22]. Let us stress however that the mechanism that allows for the increased naturalness is completely different: while in the NMSSM it is due to the enhanced tree level Higgs boson mass, here it is due to the collective loop enhancement which reduces the sensitivity to the single mass involved. A second point which is worth mentioning is that stop masses in the TeV range do not increase the fine-tuning, which is basically driven by  $m_{R_d}^2$ , Eq. (4.2). Together with the already mentioned improved naturalness bound on the gluino mass, this makes less worrisome the non observation so far of any superpartner at the LHC.

What can we expect to observe at LHC-13, given this framework? It is of course quite difficult to make a complete a solid statement. As we have seen, since the fine-



tuning is driven by  $m_{R_d}$  and  $m_{adj}$  in the interesting part of the parameter space, naturalness does not require the stop to be as light as possible. On the contrary, a relatively heavy stop (with a mass around 1 TeV) is preferred since it can give a sizable contribution to the Higgs mass, allowing for the state which are driving the fine-tuning to be lighter. In any case, we still expect  $\mu$  to be as low as possible, with the Higgsino possibly “right around the corner”.

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## A Potential minimization and mass matrices

We collect here useful formulas obtained from the minimization of the tree level scalar potential. For simplicity, we will take from the beginning the limit  $\tan \beta \gg 1$ .

Using for the vacuum expectation values the convention  $\langle h_u \rangle = v$ ,  $\langle t \rangle = v_T$ ,  $\langle s \rangle = v_S$ , the minimization of the scalar potential, Eq. (2.8) gives

$$\begin{aligned} m_{H_u}^2 &= \sqrt{2} (gM_{\tilde{W}}v_T - g'M_{\tilde{B}}v_S) - 2m_Z^2 - (\lambda_S v_S + \lambda_T v_T + \mu)^2, \\ v_T &= \frac{\sqrt{2}gM_{\tilde{W}} - 2\lambda_S\lambda_T v_S - 2\lambda_T\mu}{2(2B_T + 4M_{\tilde{W}}^2 + m_T^2 + v^2\lambda_T^2)} v^2, \\ v_S &= -\frac{\sqrt{2}g'M_{\tilde{B}} + 2\lambda_S\lambda_T v_T + 2\lambda_S\mu}{2(2B_S + 4M_{\tilde{B}}^2 + m_S^2 + v^2\lambda_S^2)} v^2. \end{aligned} \tag{A.1}$$

The squared mass matrix for the CP-even scalars, in the  $(h_u, t, s, r_d)$  basis, reads

$$M_{\text{CP-even}}^2 = \begin{pmatrix} m_Z^2 & \cdot & \cdot & \cdot \\ v(\lambda_T(\lambda_S v_S + \lambda_T v_T + \mu) - \sqrt{2}gM_{\tilde{W}}) & m_{T_R}^2 + \lambda_T^2 v^2 & \cdot & \cdot \\ v(\lambda_S(\lambda_S v_S + \lambda_T v_T + \mu) + \sqrt{2}g'M_{\tilde{B}}) & \lambda_S\lambda_T v^2 & m_{S_R}^2 + \lambda_S^2 v^2 & \cdot \\ 0 & 0 & 0 & m_H^2 \end{pmatrix}, \tag{A.2}$$

where  $m_{T_R}^2$  and  $m_{S_R}^2$  are defined below Eq. (2.11), while  $m_H^2$ , the mass of the CP-even inert doublet, is given by

$$m_H^2 = \mu^2 + m_{R_d}^2 - \frac{m_Z^2}{2} + \sqrt{2}(gM_{\tilde{W}}v_T - g'M_{\tilde{B}}v_S) + 2(\lambda_S v_S + \lambda_T v_T)\mu + (\lambda_S v_S + \lambda_T v_T)^2 + (\lambda_S^2 + \lambda_T^2)v^2 \quad (\text{A.3})$$

Turning to the CP-odd squared mass matrix, in the  $(a_t, a_s, a_{r_d})$  basis it is

$$M_{\text{CP-odd}}^2 = \begin{pmatrix} m_T^2 - 2B_T + \lambda_T^2 v^2 & \cdot & \cdot \\ \lambda_S \lambda_T v^2 & m_S^2 - 2B_S + \lambda_S^2 v^2 & \cdot \\ 0 & 0 & m_H^2 \end{pmatrix}, \quad (\text{A.4})$$

with the CP-odd component of the inert doublet degenerate in mass with the CP-even part.

To conclude, the entries of the charged scalar squared mass matrix in the basis  $(H_u^+, T^+, (T^-)^*, (R_d^-)^*)$  are

$$\begin{aligned} M_{11}^2 &= 2v_T \left[ \sqrt{2}gM_{\tilde{W}} - 2\lambda_T(\lambda_S v_S + \mu) \right] \\ M_{12}^2 &= -\frac{v}{2} \left[ \sqrt{2}g^2 v_T - 2gM_{\tilde{W}} + 2\sqrt{2}\lambda_T(\lambda_S v_S - \lambda_T v_T + \sqrt{2}\mu) \right] \\ M_{13}^2 &= \frac{v}{2} \left[ \sqrt{2}g^2 v_T + 2gM_{\tilde{W}} - 2\sqrt{2}\lambda_T(\lambda_S v_S + \lambda_T v_T + \sqrt{2}\mu) \right] \\ M_{22}^2 &= m_T^2 + 2M_{\tilde{W}}^2 + 2\lambda_T^2 v^2 + \frac{g^2}{2}(2v_T^2 - v^2) \\ M_{23}^2 &= 2(M_{\tilde{W}}^2 + B_T) - g^2 v^2 \\ M_{33}^2 &= m_T^2 + 2M_{\tilde{W}}^2 + \frac{g^2}{2}(2v_T^2 + v^2) \\ M_{44}^2 &= m_H^2 + m_{\tilde{W}}^2 - 2\sqrt{2}gM_{\tilde{W}}v_T - 4\lambda_S \lambda_T v_S v_T - 4\sqrt{2}\lambda_T \mu v_T + (\lambda_T^2 - \lambda_S^2)v^2 \end{aligned} \quad (\text{A.5})$$

with all the other entries vanishing. The  $3 \times 3$  submatrix obtained by taking out the  $R_d^-$  entry has vanishing determinant as expected, since one combination of the charged scalars is the would-be Goldstone boson eaten up by the  $W^\pm$ .

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